

SN1987A—a Testing Ground for the KARMEN Anomaly

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Abstract

We show, that SN1987A can serve as an astrophysical laboratory for testing the viability of the assertion that a new massive neutral fermion is implied by the KARMEN data. We show that a wide range of the parameters characterizing the proposed particle is ruled out by the above constraints making this interpretation very unlikely.

1 Introduction

The KARMEN collaboration, which studies neutrino from pion decay at the Rutherford laboratory has reported an anomaly in the time structure of their signal, which deviates from that expected from the standard decay into muons [1], [2]. One possible interpretation of this, offered in the experimental papers, is the production of a heavy neutral fermion with a mass of approximately 33.9 MeV and a specific range of values for the life times. It is the goal of this paper to study the (astrophysical) viability of such a fermion using supernova 1987A data.

As is well known, the KARMEN detector is located 17.5 m away from the target in the ISIS spallation neutron source. Two pulses of 800 MeV protons, of 100s duration, and $0.32\mu s$ separation, impinge every 20 milliseconds on the target [1], [2]. The time distribution of $\sim 5 \times 10^3$ events, where energy in excess of 10 MeV has been deposited in the detector, fits very well the assumption that these are interactions of neutrinos from stopped π^+ (and subsequent μ^+) decays with a moderate, time independent background. However, the time bin between 3.1 to 4.1 microseconds has an excess of ~ 100 events over of the 550 expected from the overall fit.[3]

If this 4.5σ effect is not an accidental fluctuation or an instrumental artifact, its explanation requires a slow ($\beta \sim \frac{1}{60}$) relatively monochromatic “messenger” traversing the distance $L = 17.5\text{m}$ in a time of $L/(\beta c) \sim 3.6\mu s$. It should then deposit an energy $E \geq 10\text{MeV}$ in the detector. Since many (about 10^{24}) neutrons were produced over the time, is it possible to explain the anomaly via a neutron messenger? The order of 10 meter iron shielding corresponds to $n\sigma \sim 60$ mean free paths of neutron interactions. One would then expect the neutrons to be absorbed/ diffused and in any event *not* to come in a well defined time.

However, the transport of neutrons from target to detector can be dominated by *one specific* crack or external path involving reflections. Even in this case the “neutron = messenger” hypothesis may not be viable. Indeed,

(i) The neutrons emerge from the target with high ($\sim 100\text{MeV}$) energies. There is no obvious, early moderator which is required to slow the neutrons down to $\sim 100\text{KeV}$ energies. Further, there is no obvious mechanism for guaranteeing the required time delay

$$\frac{L_{path}}{c} \left\langle \frac{1}{\beta} \right\rangle \sim 3.6\mu s \quad (1)$$

with angular brackets indicating average (over many collisions with possible moderation) along the crack/external path length, L_{path} .

(ii) The interactions in the anomalous time-bins are distributed uniformly over the detector rather than near the front edge or the termination of a putative crack.

(iii) The neutron shielding around the detector was enhanced in the second phase of the experiment. This did not stop the steady buildup of the excess of events in the anomalous bin.

(iv) The strongest argument against the neutron hypothesis comes from the energy distribution of the events. Barring the unlikely possibility that the detector had substantial amounts of ^3He and/or ^{235}U , the generic (n, γ) reactions of the slow neutrons *cannot* deposit energy in excess of the 8 MeV nuclear binding. If the observed time-anomaly is enhanced/weakened by an energy cut $E \geq 15\text{MeV}$, the case against a neutron messenger will be considerably strengthened/weakened. Unfortunately, because of possible excess of events in the first energy bin $10 \div 15$ MeV in figure 4 of [3] one cannot make this assertion yet.

For the rest of our discussion we will assume that the KARMEN anomaly has no simple explanation. We thus have to address the bold hypothesis [1, 3] that a new neutral fermion n^0 exists. The properties of this, n^0 , (the “Karmino” implied by the KARMEN data) are:

(i) Its mass is precisely tuned to be

$$m_{n^0} = m_{\pi^+} - m_{\mu^+} - 5\text{KeV} = 33.906 \pm 0.005\text{MeV} \quad (2)$$

We assume that in a (small) fraction ($= B_r(\pi^+ \rightarrow n^0 + \mu^+)$) of the cases, the stopped π^+ decays into $n^0 + \mu^+$. The monochromatic n^0 , tuned to have $\beta = 1/60$, will then arrive at the KARMEN detector $3.6\mu\text{s}$ later.

If some of n^0 particles decay via $n^0 \rightarrow e^+e^-\nu$ or via $n^0 \rightarrow \gamma\nu$ while traversing the detector, an energy $\geq 17\text{MeV}$ (or exactly 17 MeV, respectively) will be deposited. This would explain the anomaly if the $B_r(\pi^+ \rightarrow \mu^+n^0)$ and the decay rate say $\Gamma_{n^0 \rightarrow e^+e^-\nu}$ satisfy

(ii)

$$B_r(\pi^+ \rightarrow \mu^+n^0)\Gamma(n^0 \rightarrow e^+e^-\nu) = \frac{B_r(\pi^+ \rightarrow \mu^+n^0)}{\tau(n^0 \rightarrow e^+e^-\nu)} = 2.6 \times 10^{-11}\text{s}^{-1} \quad (3)$$

Is such a particle defined by equations (2), and (3) above (and most likely having weak interactions in matter) consistent with other terrestrial and astrophysical data? This issue has been considered before [4]. Recent new experimental bounds [7], and new astrophysical considerations presented below motivate us to reconsider it.

2 Terrestrial data and particle physics considerations

The particle physics constraints on the properties of the singlet fermion implied by the Karmen anomaly (Karmino) have been extensively studied in several recent papers [4, 5, 6]. Here we revisit them to set the stage for our ensuing discussion and also to incorporate the constraints from the PSI search for the Karmino[7]. We also find further constraints on the parameters of Karmino within the context of some plausible theoretical assumptions.

The striking kinematics of the $\pi^+ \rightarrow \mu^+ n^0$ decay at rest, are equally striking in decays in flight. Energetic pions yield decay muons moving in the same direction and with the same speed. This feature has been used in a recent experiment at PSI—which was inspired by the KARMEN anomaly—to obtain the remarkable bound [7]

$$B_r(\pi^+ \rightarrow \mu^+ + n^0) \leq 2.6 \times 10^{-8} . \quad (4)$$

Equation (3) then implies that

$$\tau_{(n^0 \rightarrow e^+ e^- \nu)} \lesssim 10^3 s . \quad (5)$$

The lack of evidence for sharp ~ 17 MeV energy deposition for events in the anomalous bin, and a theoretical bias (the need to use loops rather than tree diagram) suggest that the above decay, rather than $n^0 \rightarrow \nu + \gamma$, dominates.

It is useful to parameterize the $\pi \rightarrow \mu^+ n^0$ and $n^0 \rightarrow e^+ e^- \nu$ processes via effective local four-Fermi interactions.

$$\tilde{G}_\mu \bar{\Psi}_d(x) \gamma_5 \Psi_u(x) \Psi_{n^0}(x) \Gamma' \Psi_\mu(x) , \quad (6)$$

and

$$\tilde{G}_{n^0} \bar{\Psi}_{n^0}(x) \Gamma'' \Psi_{\nu_d}(x) \Psi_e(x) \Gamma' \Psi_e(x) , \quad (7)$$

with Lorentz/flavor structure which may be different from those for neutrinos in the standard Electroweak model. Still we can compare using just phase space considerations¹, the expected ratio of the decay rates: $\Gamma_n/\Gamma_\mu \approx (m_{n^0}/m_\mu)^5 \tilde{G}_\mu^2/G_F^2$. Using $m_{n^0}/m_\mu \approx 1/3$, $\Gamma_\mu \approx 4 \times 10^5 s^{-1}$ and equation (5) we find that

$$\frac{\tilde{G}_{n^0}^2}{G_F^2} \gtrsim 0.7 \times 10^{-6} . \quad (8)$$

The local effective Lagrangian implies also the crossed scattering process (at energies $E_\nu \gg m_{n^0}^2/2m_e \sim \text{GeV}$)

$$\bar{\nu}_d + e^- \rightarrow e^- + n^0 , \quad (9)$$

with ν_d the neutrino appearing in the n^0 decay $n^0 \rightarrow e^+ e^- \nu_d$, and the process $e^+ e^- \rightarrow n^0 \bar{\nu}_d$ with crosssection

$$\frac{\sigma(e^+ e^- \rightarrow n^0 \bar{\nu}_d)}{\sigma(e^+ e^- \rightarrow \nu_i \bar{\nu}_i)} \sim \frac{\tilde{G}_{n^0}^2}{G_F^2} . \quad (10)$$

This last process features in SN1987A and supernovae in general. We therefore need to know the constraints on the coupling \tilde{G}_{n^0} from laboratory experiments to study the astrophysical viability of the Karmino.

The specific Lorenz structure of the standard $\pi^+ \rightarrow \mu^+ \nu$ (or $\pi^+ \rightarrow n^0 \mu^+$) decay only mildly affects the decay rate. Hence,

$$\frac{\tilde{G}_\mu^2}{G_F^2} \beta \approx \frac{\Gamma(\pi^+ \rightarrow n^0 \mu^+)}{\Gamma(\pi^+)} \equiv B_r(\pi^+ \rightarrow n^0 \mu^+) \leq 2.6 \times 10^{-8} , \quad (11)$$

¹For more detailed discussion that takes into account the general Lorentz structure carefully, see [8].

where we included the phase space suppression factor $\beta \sim 1/60$ for $\pi^+ \rightarrow n^0 \mu^+$. From the last equation we find that

$$\frac{\tilde{G}_\mu^2}{G_F^2} \lesssim 1.5 \times 10^{-6} . \quad (12)$$

Very short partial life-time in the $10^{-6}\text{s} \leq \tau(n^0 \rightarrow e^+ e^- \nu_d) \leq 0.6 \times 10^{-3}\text{s}$ range can be excluded by the following “theoretical arguments” showing that no new interactions stronger than the ordinary “weak” interactions are allowed. Indeed, $\tilde{G}_{n^0} = G_F \sqrt{\epsilon_{n^0}}$ implies that

$$\tau(n^0 \rightarrow e^+ e^- \nu_d) = \tau(\mu) \left(\frac{m_\mu}{m_{n^0}} \right)^5 \epsilon_{n^0}^{-1} = 0.7 \times 10^{-3} \epsilon_{n^0}^{-1} \text{s} . \quad (13)$$

From this we see that life-times shorter than 10^{-3} seconds would require values of \tilde{G}_{n^0} larger than G_F . However, in the presently accepted approach to field theories one *cannot* postulate at will new non-renormalizable four-Fermi interactions. The latter should be viewed only as the low energy limit of massive boson exchanges. The bilinear vertices of the boson exchange tree-diagram, $\bar{\Psi} B \Psi$, which are parts of the interaction Lagrangian in a fully renormalizable and acceptable field theory. To date, the only such known bosons are the vectorial (spin 1) gluons, photon, and W^+ , W^- , Z^0 which generate the strong, electromagnetic and weak interactions, respectively. Yet, the exchanges could also be scalars generating $aS + bP$ four vertices instead of the (V-A) combination for W_μ .

The effective \tilde{G}_{n^0} is then of the form $\tilde{g}_1 \tilde{g}_2 / m_{\tilde{X}^2} \propto \tilde{G}_{n^0}$, with $\tilde{g}_1 \tilde{g}_2 v$ representing the $\tilde{X} e^+ e^-$ and $\tilde{X} n^0 \nu_d$ vertices and $m_{\tilde{X}}$ is the mass of the new boson. This should be compared with $g_W^2 / m_W^2 \propto G_F$ (with the same proportionality constant). Since, $g_W \approx 0.7$ and we do not wish to entertain $\tilde{g} > 1$, we need $m_X \leq m_W$ in order to allow for $\tilde{G} > G_F$. However, we will now have either a charged X^- particle exchange in $e^- \nu_d \rightarrow X^- \rightarrow e^- n^0$ or (and) a neutral X^0 in the crossed channel.: $e^+ e^- \rightarrow X^0 \rightarrow \bar{\nu}_d n^0$, with m_{X^-} or (and) m_{X^0} smaller than m_W .

The new LEP runs at $W_{CM} \geq 200 \text{ GeV}$ strongly exclude new charged particles of mass smaller 85 GeV. Also, the new X^0 would have strongly manifested as a new narrow resonance in electron-positron scattering in the Tristan and LEP regimes. (We note also that for electron-positron scattering at energies larger than the mass of X the latter could still manifest if there is an emission of an initial internal bremsstrahlung photon.)

The LEP bounds ($N_\nu \lesssim 3.01$) on sequential, $I_{weak} \neq 0$, particles with masses $\leq 40 \text{ GeV}$ imply that n^0 is a weak-Iso-spin singlet. If the decays $\pi^+ \rightarrow \mu^+ n^0$ and/or $n^0 \rightarrow e + e^- \nu_d$ do involve also left-handed quarks or leptons then by using $SU(2)_W$ rotation the processes could be related to their isospin analogues. Thus $\pi^+ \rightarrow \mu^+ n^0$ would be related to $\pi^0 \rightarrow \bar{\nu} n^0$ which is controlled by the same \tilde{G}_μ . The latter may again manifest via $\nu_\mu + N \rightarrow N' + n^0$ in the hot proto neutron star. The latter process is more effective there (has a smaller Boltzman factor suppression) than $\mu^+ n \rightarrow p + n^0$. Note, however that the charged current process is guaranteed by the very existence of four-Fermi coupling, equation (6), with no additional assumptions.

The $\pi^+ \rightarrow \mu^+ n^0$ process suggests that, if muonic lepton number is conserved then n^0 has $L_\mu = 1$ and hence the decay neutrino ν_d in $n^0 \rightarrow e^+ e^- \nu_d$ should also have $L_\mu = 1$. If we make the natural (though not mandatory!) assumption that $\nu_d = \bar{\nu}_\mu$, then by crossing and isospin rotation, we could generate yet another decay mode for the muon.

$$\mu^+ \rightarrow e^+ \nu_e n^0 . \quad (14)$$

Apart from a small ($\sim 40\%$) phase space correction due to $m_{n^0}/m_{\mu^+} \sim 1/3$ the rate of the new mode is given just like for $\Gamma(\mu \rightarrow e\nu\nu)$ by

$$\Gamma(\mu^+ \rightarrow e^+ \nu_e n^0) \approx \frac{\bar{G}_\mu^2 m_\mu^5}{192\pi^3}, \quad (15)$$

hence

$$\frac{\Gamma(\mu^+ \rightarrow e^+ \nu_e n^0)}{\Gamma(\mu^+ \rightarrow e^+ \nu \nu)} \approx \frac{\bar{G}_\mu^2}{G_F^2} \equiv \epsilon_\mu. \quad (16)$$

Since the muon decay rate and the electron spectrum and polarization in this decay have been studied carefully over the last four decades, one could deduce that :

$$\epsilon_\mu \leq 10^{-3} \left(\text{if } n^0 \rightarrow e^+ e^- \nu_\mu \right) \quad (17)$$

Using equation (8) this implies a lower bound on $\tau_{n^0 \rightarrow e^+ e^- \nu_\mu}$:

$$\tau_{n^0 \rightarrow e^+ e^- \nu_\mu} \geq 1\text{s} \quad \left(\text{if } n^0 \rightarrow e^+ e^- \nu_\mu \right), \quad (18)$$

in addition to the model-Independent upper bound $\tau_{n^0 \rightarrow e^+ e^- \bar{\nu}_d} \leq 10^3$ s implied by the PSI upper bound on $B_r(\pi \rightarrow \mu^+ n^0)$.

To summarize, we note that the PSI experiments in conjunction with some plausible theoretical assumptions lead to bounds on the lifetime of n^0 : $1\text{ s} \leq \tau_{n^0} \leq 10^3\text{ s}$. Furthermore, the strength of the four-Fermi coupling of n^0 to electrons defined by the parameter ϵ_{n^0} (see the previous section) has a lower bound $\epsilon_{n^0} \geq 0.7 \times 10^{-6}$ which will imply lower bounds on the production of n^0 in supernovae environments. We consider the impact of this on SN1987A observations in the next section.

3 Limits Derived from Supernovae, Notably Supernova 1987A

In the standard scenario for a type-II supernova [9], most ($\sim 99.7\%$) of the collapse energy is emitted via neutrinos. The latter are trapped in the dense hot proto neutron star for few seconds during which they diffuse out. The optical signal is delayed by $\sim 1 \div 3$ days, after the shock emanating from the core traverses the progenitor's envelope. The energy of the shock ($\sim 10^{51}$ erg) manifest as the kinetic energy of the ejected envelope. Only about a percent i.e, 10^{49} erg manifest in the (bolometric) electromagnetic signal.

This scenario— largely worked out before the SN87A— explosion has been brilliantly verified there. A notable exception is the fact that the progenitor was not a red supergiant but a blue supergiant. Accordingly, the optical signal appeared an usually short time of ~ 3 hours after the neutrino pulse that marked the core collapse.

Many putative novel, features in particle physics can modify the above scenario and/or leave other enduring traces. Hence, considerations of SN1987A data (and supernovae in general) strongly limited [10] the interactions of axions, right handed neutrinos, (generated via m_{Dirac} and/or via precession of neutrino spin in magnetic field when the neutrino posses an anomalous magnetic moment [11], [12]). Finally electromagnetic neutrino decays are severely limited.

We would like to find the ranges of lifetime τ and branching B_r allowed by the observations from SN1987A. For most part we will only assume that $B_r/\tau \sim 2.6 \times 10^{-11} \text{s}^{-1}$ as required by the magnitude of the KARMEN anomaly. In terms of the two four-Fermi couplings introduced above this equivalent to

$$\epsilon_{n^0} \epsilon_\mu \equiv \frac{\tilde{G}_{n^0}^2}{G_F^2} \frac{\tilde{G}_\mu^2}{G_F^2} \approx 10^{-12} . \quad (19)$$

We note that $B_r \approx 2.6 \times 10^{-8}$ (the upper bound of the recent PSI experiment) and the corresponding $\tau \approx 10^3 \text{s}$ imply roughly equal ϵ_{n^0} , ϵ_μ (actually $\epsilon_{n^0} = 0.6 \times 10^{-6}$ and $\epsilon_\mu = 1.5 \times 10^{-6}$). As will be shown below, the existence of the n^0 s with such properties can have quite dramatic effects on supernovae astrophysics.

3.1 Neutrino emission

The production of n^0 in the newly formed hot core can proceed via both the \tilde{G}_{n^0} and \tilde{G}_μ couplings. That is via

$$e^+ e^- \rightarrow n^0 + \nu_d, \quad (20)$$

with crossection

$$\sigma \approx \tilde{G}_{n^0}^2 E_e^2 , \quad (21)$$

and via nuclear scattering

$$\nu N \rightarrow n^0 N , \quad (22)$$

with crossection

$$\sigma \approx \tilde{G}_\mu^2 E_\nu^2 . \quad (23)$$

The latter process can be viewed as due to π^0 exchange. At low (≤ 50 MeV) energies this yields an effective local four-Fermi vertex. If the $\pi^0 \rightarrow \nu n^0$ vertex *cannot* be inferred from $\pi^+ \rightarrow \mu^+ n^0$ we have instead $\mu^+ n \rightarrow n^0 p$ with crossection

$$\sigma(\mu^+ n \rightarrow n^0 p) \approx \tilde{G}_\mu^2 m_\mu^2 . \quad (22')$$

These processes should be compared with the standard neutrino scattering crossection, $\sigma_W \approx G_F^2 E_\nu^2$. In the dense ($\rho \gtrsim 10^{15} \text{gcm}^{-3}$, i.e. super-nuclear density) and hot ($T \gtrsim 30 \text{MeV}$) core, the resulting neutrino mean free path is:

$$l = \frac{1}{n \sigma_W} = 100 \text{cm}. \quad (24)$$

It implies a total number of neutrino collisions, during diffusion

$$N_{\nu N \text{ collisions}} \approx \left(\frac{R}{l} \right)^2 \gtrsim 10^8, \quad (25)$$

where $R \approx 30 \text{km}$ is adopted for the radius of the very hot core. In a fraction $\epsilon_\mu f_B(m_{n^0}/T)$ of the collisions we could have, instead of the usual weak scattering, n^0 production $\nu_\mu + N \rightarrow n^0 + N$, or $\mu^+ n \rightarrow p + n^0$.

Note that the Boltzman factor, f_B , reflects the energy distribution of the initial particles (neutrinos in the present case). The n^0 themselves need not be in thermal equilibrium. For $T \sim 30$ MeV the appropriate $f_B \approx (m/T)^{3/2} e^{-m_\mu/T}$ is ~ 1 and for

$T \sim 10$ MeV it is ~ 0.5 . When only $\mu^+ n \rightarrow p + n^0$ contributes, the corresponding Boltzman factor is $f'_B \approx 2.7 (E/T)^{3/2} e^{-m_\mu/T}$ which for $T = 10$ MeV is $\sim 1/4$.

When $\epsilon_{n^0} \gg \epsilon_\mu$ n^0 production is dominated by $e^+ e^- \rightarrow n^0 \nu_d$. To estimate the rate of the latter, we note that all collapse models indicate a sizable electron density in the proto-neutron star ($\frac{n_e}{n_N} \geq 0.2$ [14]) during the first few seconds. The number of weak interactions induced $e^+ e^-$ collisions per positron during these few seconds is now given by:

$$N_{Weak-collisions} \approx \left(\frac{R}{l_W} \right)^2 \approx 10^6, \quad (26)$$

where we have used $l_W = (n_e \sigma_W)^{-1} \approx 1000$ cm.

Since $\frac{\sigma(\nu_\mu + N \rightarrow n^0 + N)}{\sigma(\nu_\mu + N)} \approx \frac{\tilde{G}_\mu^2}{G_F^2} \approx \epsilon_\mu$, and $\frac{\sigma(e^+ e^- \rightarrow n^0 + \nu_d)}{\sigma(e^+ e^- \rightarrow \nu \bar{\nu})} \approx \frac{\tilde{G}_{n^0}^2}{G_F^2} \approx \epsilon_{n^0}$, n^0 production occurs in a fraction of $\epsilon_\mu f_B$ (or $\epsilon_{n^0} f_B$) of the $10^8 \nu N$ (or the 10^6 weak $e^+ e^-$) collisions. Altogether, the number of produced n^0 s is $N_{total}(n^0) \approx N_{total}(\nu_i) (10^8 \epsilon_\mu + 10^6 \epsilon_{n^0}) f_B \geq 20 f_B N_{total}(\nu_i) \approx (20 \div 5) N_{total}(\nu_i)$. In the penultimate step we used $x + y \geq 2\sqrt{xy}$ and $\epsilon_\mu \epsilon_{n^0} = 10^{-12}$, and finally substituted $f_B = 1$ (or 0.6) for reactions r_1, r_2 respectively and $T = 10$ MeV.²

Note the the above rate of n^0 particle production (which already tends to exceed by $5 \div 20$ the initial number of parent neutrinos) is actually a minimum achieved for a particular choice of $\epsilon_\mu / \epsilon_{n^0} \approx N_{\nu N \text{ collisions}} / N_{e^+ e^- \text{ Weak collisions}} \approx 100$. Generically equation (29) above leads to $N_{n^0} > f N_{\nu_i}$ with $f \sim 10^2 \div 10^6$! n^0 particles³. What this apparently paradoxical result, really means is the following. Our assumption that neutrinos continuously convert into n^0 s during their entire ordinary, diffusion time $t_{diff} \approx l_{meanfreepath} N_{coll} / c \sim 0.3$ s is false. Rather, a sizeable fraction converts into n^0 on a much shorter timescale $t_{convers} \approx t_{diff} / f \ll 0.1$ s.

If the n^0 particles escape on this short timescale, out of the hot proto-neutron star, the momentary reaction balance will immediately shift toward the n^0 s. Catastrophic cooling of the core on this short timescale will then ensue. The resulting neutrino pulse would be drastically shortened in time and reduced in intensity. The IMB and Kamiokande observation of SN1987A neutrino pulses lasted $5 \div 10$ seconds and indicated that the total energy emitted via neutrinos was $\gtrsim 3 \times 10^{53}$ erg, as expected. Hence, these observations by themselves exclude almost the complete range of B_r, τ (or $\epsilon_\mu, \epsilon_{n^0}$) parameters and thereby the n^0 hypothesis itself.

To complete the argument let us estimate the escape time from the core for any single n^0 particle. The n^0 particles diffuse out on timescale of $t_{n^0, diff} \ll t_{diff} / 100 \approx 3$ ms, once the crosssection of the n^0 s on ambient particles ($\sigma_{n^0 N}, \sigma_{n^0 e}, \sigma_{n^0 \nu}$) is ten times smaller

² Due to the chemical potential in the supernova core there is a suppression of the positron density in the supernova core, which is not explicit in our discussion, since the effect cancels in $\frac{\sigma(e^+ e^- \rightarrow n^0 + \nu_d)}{\sigma(e^+ e^- \rightarrow \nu \bar{\nu})} \approx \frac{\tilde{G}_{n^0}^2}{G_F^2} \approx \epsilon_{n^0}$.

The suppression factor is $\sim \left(\frac{T}{\mu} \right)^3 e^{-\mu/T}$. Taking $\mu \sim 200$ MeV and $T \sim 60$ MeV, this factor is of order 2×10^{-3} . There is a more dominant mechanism for n^0 production that comes from $\nu_\mu + e^- \rightarrow e^- + n^0$ which does not suffer from chemical potential suppression. As long as the muon neutrinos are equally abundant in the core, this process adds to the first term in $N_{total}(n^0) \approx N_{total}(\nu_i) (10^8 \epsilon_\mu + 10^6 \epsilon_{n^0}) f_B \geq 20 f_B N_{total}(\nu_i) \approx (20 \div 5) N_{total}(\nu_i)$, and leaves our conclusions unchanged.

³ If only (22') rather than (20) operates then $10^6 \epsilon_{n^0} \rightarrow 10^6 \epsilon_{n^0} (f'_B / f_B)$ but this decrease the total number of the n^0 particles by merely a factor of 5.

than the corresponding σ_W (the neutrino or other particles weak crosssection). We will therefore make the assumption:

A1: The nuclear and other crosssections of n^0 particles are (at least ten times) smaller than those of neutrinos.

This assumption is clearly consistent with the value of B_r and τ (or \tilde{G}_μ^2 and \tilde{G}_μ^2) required for the KARMEN anomaly. The n^0 particles are devoid of color. Otherwise they would Hadronize and certainly could not penetrate 10 meters of iron. They are electromagnetically neutral and as noted above are also singlets of the full $SU(2) \times U(1)$ electroweak gauge group. Hence assumption A1 is eminently reasonable, though given the meager information on n^0 particles it is not however, mandated in an absolutely model-independent manner.

In all the above discussions, we made the implicit assumption that τ_{n^0} the actual decay life-time of n^0 is longer than $R/c \sim 10^{-4}$ s. Therefore the bounds derived apply only for τ values exceeding $\sim 10^{-4}$.

Thus, n^0 with lifetime in the range $1 \text{ s} \leq \tau_{n^0} \leq 10^3 \text{ s}$ would be inconsistent with the SN1987A observations. A simple way to see this is to note that the corresponding bounds for ϵ_{n^0} i.e. $10^{-6} \leq \epsilon_{n^0} \leq 10^{-3}$ are in the range forbidden for general coupling of weakly interacting particles to electrons[13].

One way to avoid this bound would be to give up the theoretically plausible assumption that n^0 has muonic quantum number and $\epsilon_{n^0} \sim 1$ in which case all the produced n^0 's get trapped and this SN1987A constraint can be avoided. In this case, both the n^0 and ν_d have to be sterile neutrinos. In this case also one can derive a constraint on ϵ_{n^0} from big bang nucleosynthesis. The point is that if ϵ_{n^0} is indeed of order one, it will inject sterile neutrinos ν_d into the cosmic soup and they will count as one neutrino species. Thus if the BBN constraint on additional neutrino species is much less than one [16], then the n^0 lifetime has to be less than 10^{-5} sec so that any population of ν_d 's injected into the cosmic soup by n^0 decays would be diluted by the QCD phase transition. This implies roughly that $\epsilon_{n^0} \geq 0.5$.

Another possibility is to postulate the existence of exotic interaction of both the n^0 and the ν_d to quarks so that those interactions trap the n^0 's in the supernova core. Clearly, such interactions must have a very specific property [17] i.e. the n^0 and ν_d must couple only to the isosinglet current so that the π^+ 's do not decay directly into the n^0 in order to be consistent with laboratory observations i.e. π^+ decay mostly to the $\mu^+ n_0$. As shown in [17], this also helps to avoid BBN constraints. To be a bit more quantitative, we may note that if the interaction is denoted by $\mathcal{L}_{int} = \frac{G_F \epsilon}{\sqrt{2}} \bar{n}^0 \Gamma \nu_d (\bar{u} \Gamma u + \bar{d} \Gamma d)$, then the n^0 -nucleon scattering cross section is given by $\sigma_{n^0 N} \sim \frac{9}{\pi} (g_V^2 + 3g_A^2) G_F^2 \epsilon^2 E^2$. For an average n^0 energy of 100 MeV, this leads to a mean free path $l_{n^0} \sim 1.5/\epsilon^2$ so that trapping in the supernova can occur for $\epsilon \geq 4 \times 10^{-3} - 3 \times 10^{-2}$. As has been noted in [17], this value of ϵ is quite adequate to suppress the contribution of n^0 to nucleosynthesis.

It should however be noted that while this can evade the energy constraints, they will be of no help in eliminating the constraint from considerations of shock energy and propagation time to be discussed in the last section of the paper.

3.2 Galactic e^+e^- population and Early γ Ray Flash from SN1987A

As indicated above the observed neutrino signal from SN1987A provides a sweeping argument against the n^0 hypothesis. Here we note that the timing, intensity and character of the ordinary electromagnetic signals from SN1987A, and from other supernovae, also exclude a wide range of τ (or equivalently B_r) values. Moreover, commulative effects of all past supernovae in the history of the galaxy also lead to strong constraints. The key assumption required to enable these constraints is :

A2: *The electromagnetic decay mode $n^0 \rightarrow e^+e^-\nu_d$ is (one of) the dominant decays of n^0 .*

Note that A2 fails to hold if invisible decays like $n^0 \rightarrow 3\nu$ dominate. This assumption implies that $\tau(n^0 \rightarrow e^+e^-\nu_d)$ is also (close to) the actual τ_{n^0} , the lifetime of n^0 .

For $\tau \gtrsim 100$ s at least 20% of the n^0 emitted will decay outside the progenitor of SN1987A radius ($\sim 3 \times 10^{12}$ cm) where the average velocity of n^0 was assumed to be $\beta \sim 0.6$. The total number of decay positrons is then $N_{e^+} \approx N(n^0 \text{ decays for } r > R) \approx 0.3N_{n^0} \approx 0.2f_B N_{\nu_i} \approx 4 \times 10^{54} \div 2 \times 10^{56}$, where $N_{\nu_i} \approx 10^{57}$ is the total number of emitted neutrinos of any given species (say $\bar{\nu}_\mu$). If the n^0 production is dominated by $\mu^+ + n \rightarrow p + n^0$ rather than by $\bar{\nu}_\mu + N \rightarrow N' + n^0$ and we use $T \approx 10$ MeV the f_B can be as small as 10^{-3} . Following [19] we note that the decay positrons could manifest in two different ways:

(i) positrons can annihilate on electrons (from decays of different n^0 particles as well as ambient) in the vicinity of the supernova progenitor.

(ii) Positrons that manage to escape into the interstellar medium, slow down and keep accumulating over the galactic history. A steady rate of annihilation yielding the monochromatic 512 KeV, γ -Ray line is then established. Dar Goodman and Nussinov in a work that preceded the explosion of SN1987A, focused on this aspect. It was found that even for f_B as small as 10^{-3} , an extremely broad range: $\tau_\nu \geq 10^4$ s⁴ is excluded by the observational limits on the 512KeV annihilation line. The occurrence of SN1987A and the lack of any observations of an early γ -Ray pulse, delayed by $R/c \sim 100$ s ($R \approx 3 \times 10^{12}$ cm is the radius of the blue supergiant progenitor of SN198978A), with respect to the neutrino signal, strongly limit electromagnetic decay at $\tau \approx 100$ s. Indeed for $\tau \lesssim 100$ s most of the n^0 s will decay within a distance $r \approx R$ from the progenitor. The resulting electron column density $\bar{n}_e R \approx \frac{N_{n^0}}{4\pi R^2} \approx \frac{N_{e^+}}{4\pi R^2} = 10^{27} \div 10^{30} \text{cm}^{-2}$, of electrons or positrons suffices as $\bar{n}_e R \sigma_{\text{annihilation}} \gg 1$, ensuring almost complete annihilation. This results in a γ fluence $\sim \frac{N_{e^+}}{4\pi D^2} \approx 3 \times (10^5 \div 10^8) \text{cm}^{-2}$

corresponding to energy fluence $(2 \div 200) \text{erg cm}^{-2}$, for a distance of 50kpc to SN1987A. Such a fluence is $6 \div 10$ orders of magnitude larger than the observational upper bound ! [9]

Hence, once $N_{e^+}/N_{\nu_i} \gtrsim 10^{-4}$ i. e. once $f_B \geq 10^{-3}$ and $\Gamma(n^0 \rightarrow e^+e^-\nu_d)/\Gamma(n^0) \approx O(1)$ (namely A2 holds) we can exclude not only $\tau \approx 100$ s but even $\tau \approx 10$ s as well, when the number of n^0 decaying outside the progenitor is suppressed by an additional multiplier of $e^{-R/(c\beta\tau)} \approx e^{-16!}$ (when $\beta \sim 0.6$).

Next, we turn our attention to even shorter lifetimes and show that also n^0 decays within the progenitor could have dramatic observational imprints on SN1987A. The lack of these imprints, severely constraining the allowed τ values.

⁴such a large τ is required for the particles to traverse red giants envelopes

3.3 Shock Energy and Propagation time

In the standard type-II supernova scenario, the bounce shock from the collapsed neutron core propagates outwards and traverses the progenitor mantle and envelop. In the process, almost all the shock energy is converted into bulk kinetic energy of the ejected envelop. When the shock reaches the surface of the star (actually when it gets to optical depth $\lesssim 1$) the heated up stellar gas emits radiation in the UV and optical. This marks the beginning of the optical light curve. Thus, for SN19897A, the numerical hydro-dynamical models [15], [18] indeed are capable of reproducing the 3 hr time delay between the neutrino pulse and the beginning of the optical light curve. The propagation time [18] can be approximated as

$$T_{prop} = 1.4 \left(\frac{R}{R_{\odot}} \right) \left(\frac{M_{ej}}{10M_{\odot}} \right)^{1/2} \left(\frac{E}{1 \times 10^{51} \text{erg}} \right)^{-1/2} \text{hr} \quad (27)$$

where M_{ej} , the ejected mass, is the mass exterior to the radius where the energy E is deposited.

In the present case, one expects to have in addition to the standard (weakened) bounce shock, a shock resulting from the thermalisation (on spot essentially) of annihilated e^+e^- resulting from the decay of the n^0 particles to e^+e^- that occurs within the mantle or envelope. This will occur for lifetime of the n^0 particles $\tau \lesssim 30$ s and at a mean radius of $r = v_{n^0} \tau$, with a spread in r which is of order r . The energy in this shock is an order of magnitude larger than that in the standard bounce shock. The total amount of heat energy deposited by the n^0 s is actually their share in the thermal energy of the proto-neutron star core, i.e. $E_{exp} \sim 3 \times 10^{52} \text{erg}$. This is in contrast with the small fraction of the neutrino energy that is converted into shock energy. Thus one expects this shock to be both more vigorous and also arrive earlier at the stellar surface,

Assuming emission at a temperature of $5 \div 10 \text{MeV}$ $\beta_{n^0} \approx 0.7$. Using the initial stellar structure given by [18] which is based on the model of [15], In the left panel of figure 1 we present the propagation time computed from equation (27) as function of τ .

We see that for $\tau \gtrsim 0.03 \text{s}$ the propagation time is substantially shorter than the observational value of 3hr. In the present framework, a propagation time of 3hr requires the progenitor to have a much larger radius which would have rendered it a *red* supergiant while it is known to have been a blue supergiant at the onset of core collapse. Also, the fact that the hydrogen-rich envelop would have been pushed out before the (now weakened) bounce shock from the core collapse could have reached it, is expected to substantially reduce the mixing of hydrogen into the iron rich layers, contrary to the observations.

3.4 Summary and Conclusions

We have seen above that the various considerations pertaining to SN1987A, and supernovae in general, rule out a large range of possible values of τ (B_r). In the right panel of figure 1, we present for each set of considerations, the range of τ values which are ruled out. We also present in the same figure the regions of τ which are excluded by the various particle physics considerations presented in §2. Some overlap, so that certain ranges of τ are excluded by more than one argument. The wide range of the ruled out τ values strongly suggests that the n^0 hypothesis is not viable.

We dedicate this paper to our colleague Arnon Dar on the occasion of his 60th birthday. We would like to thank F. T. Avignoni, C. Rosenfeld, and S. Mishra for early discussions

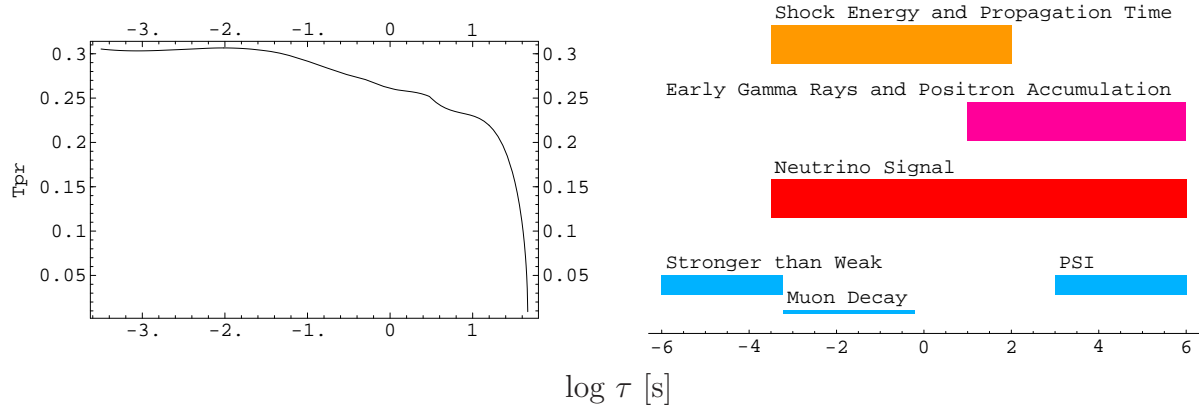


Figure 1: Left: Shock propagation time, in hours, as function of the life-time in seconds. Right: Ranges of life-times excluded by various effects

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